

Decays of the Littlest Higgs Z_H and the Onset of Strong Dynamics

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Abstract

The Little Higgs mechanism, as realized in various models, requires a set of new massive gauge bosons, some of which mix with gauge bosons of the Standard Model. For a range of mixing angles the coupling of gauge bosons to scalars can become strong, ultimately resulting in a breakdown of perturbative calculation. This phenomenon is studied in the Littlest Higgs model, where the approach to strong dynamics is characterized by increasing tree-level decay widths of the neutral Z_H boson to lighter gauge bosons plus multiple scalars. These increasing widths suggest a distinctive qualitative collider signature for the approach to the strong coupling regime of large Higgs and other scalar multiplicities. In this work we catalog the kinematically allowed three-body decays of the Z_H , and calculate the partial width of the process $Z_H \rightarrow Z_L H H$. This partial width is found to be larger than the comparable two-body decay $Z_H \rightarrow Z_L H$ for values of the $SU(2)$ mixing angle cosine $\theta < 0.13$, indicating divergence of the Littlest Higgs sigma field expansion at values of cosine θ larger than a simple parametric calculation would suggest. Additionally, we present analytical expressions for all two-body decays of the Littlest Higgs Z_H gauge boson, including the effects of all final-state masses.

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I. INTRODUCTION

A well known formal difficulty of the Standard Model of particle physics is that the Higgs boson mass has quadratically divergent contributions from fermion and gauge boson loops which must be cancelled by a ‘fine-tuned’ tree-level mass term. This fine-tuning is removed in supersymmetric models by introducing boson partners for Standard Model fermions, and fermion partners for Standard Model bosons. Cancellations of quadratic divergences then arise from the differing signs of bosonic and fermionic loops. Recently a new class ‘Little Higgs’ models[1], has emerged in which cancellations occur between paired particles of the *same* statistics: fermions are paired with new, massive fermions, and bosons with new massive bosons.

This Little Higgs mechanism, as realized in various models, requires a set of new massive gauge bosons, some of which mix with gauge bosons of the Standard Model. For some range of mixing angles the coupling of gauge bosons to scalars can become strong, resulting in a breakdown of perturbative calculation. In this work we probe the limits of perturbative reliability by examining the decay partial widths of a massive gauge boson.

The prototype Little Higgs model is the ‘Littlest Higgs’[2]. Here we study the decays of a new neutral $SU(2)$ gauge boson in the Littlest Higgs model, the Z_H . The presence and properties of this particle are central to the divergence-canceling mechanism of the Littlest Higgs model.

A catalog of Feynman rules for the Littlest Higgs model can be found in Reference [3], along with a number of phenomenological results. A few important corrections can be found in Reference [4]. Additional rules for scalar couplings can be found in Reference [5]. Here we will sketch the outlines of the model as needed to establish notation and motivate our work.

The enlarged electroweak space of the Littlest Higgs model is the symmetric tensor representation of $SU(5)$, which is broken at a scale f by a tensor vacuum expectation value to the coset $SU(5)/SO(5)$. The result is $24 - 10 = 14$ Goldstone bosons, which are parameterized by a 5×5 non-linear sigma field $\Sigma(x)$. Four of these Goldstone bosons will become the complex Higgs doublet of the Standard Model, four will become the longitudinal modes of four new massive gauge bosons, and the remaining six will form a massive complex scalar triplet.

Particles	Spin	Mass
$\Phi^0, \Phi^P, \Phi^+, \Phi^{++}, \Phi^-, \Phi^{--}$	0	$\frac{\sqrt{2}m_H f}{v(1-(\frac{4v'f}{v^2})^2)^{\frac{1}{2}}}$
$\mathbf{T}, \bar{\mathbf{T}}$	$\frac{1}{2}$	$\frac{v}{m_t}(\lambda_1 \lambda_2 f)$
\mathbf{A}_H	1	$m_Z s_w (\frac{f^2}{5s'^2 c'^2 v^2} - 1)^{\frac{1}{2}}$
\mathbf{Z}_H	1	$m_W (\frac{f^2}{s^2 c^2 v^2} - 1)^{\frac{1}{2}}$
$\mathbf{W}_H^+, \mathbf{W}_H^-$	1	$m_W (\frac{f^2}{s^2 c^2 v^2} - 1)^{\frac{1}{2}}$

TABLE I: New particles in the Littlest Higgs model [3][4], where $m_W = \frac{gv}{2}$, $m_Z = \frac{gv}{2c_w}$.

The $SU(5)$ group contains an $SU(2)_1 \otimes U(1)_1 \otimes SU(2)_2 \otimes U(1)_2$ subgroup, which is gauged, with gauge couplings g_1, g'_1, g_2 , and g'_2 , respectively. The result is eight gauge bosons, which after diagonalization to mass eigenstates are the gauge bosons of the Standard Model, here labeled A_L, Z_L, W_L^\pm , and new gauge bosons with masses of order f , labeled A_H, Z_H, W_H^\pm . The mass diagonalization can be described in terms of the mixing angles:

$$c = \cos \theta = \frac{g_1}{\sqrt{g_1^2 + g_2^2}} = \sqrt{1 - s^2} \quad (1)$$

and

$$c' = \cos \theta' = \frac{g'_1}{\sqrt{g'^2_1 + g'^2_2}} = \sqrt{1 - s'^2}. \quad (2)$$

With $SU(5)$ represented in this way, quadratic Higgs mass divergences from Standard Model gauge boson loops are cancelled by opposite-sign contributions from the new massive gauge bosons, leaving a logarithmic divergence. At some scale above 10 TeV this logarithmic contribution must also be cancelled if fine-tuning is to be avoided, so we do not view the Littlest Higgs model as complete, but instead expect at sufficiently high energies additional contributions from an unidentified complete theory, the so-called ‘‘ultraviolet completion’’ [6].

The requirement that Yukawa terms be gauge-invariant restricts the $U(1)_1$ and $U(1)_2$ ‘hypercharge’ assignments of the fermions, such that the sum for each fermion equals the Standard Model hypercharge [3]. The remaining freedom of assignment is limited to one parameter for the quarks, y_u , and one for the leptons, y_e . This freedom is eliminated if we require that all anomalies are cancelled in the theory, resulting in the fixed values $y_u = -\frac{2}{5}$

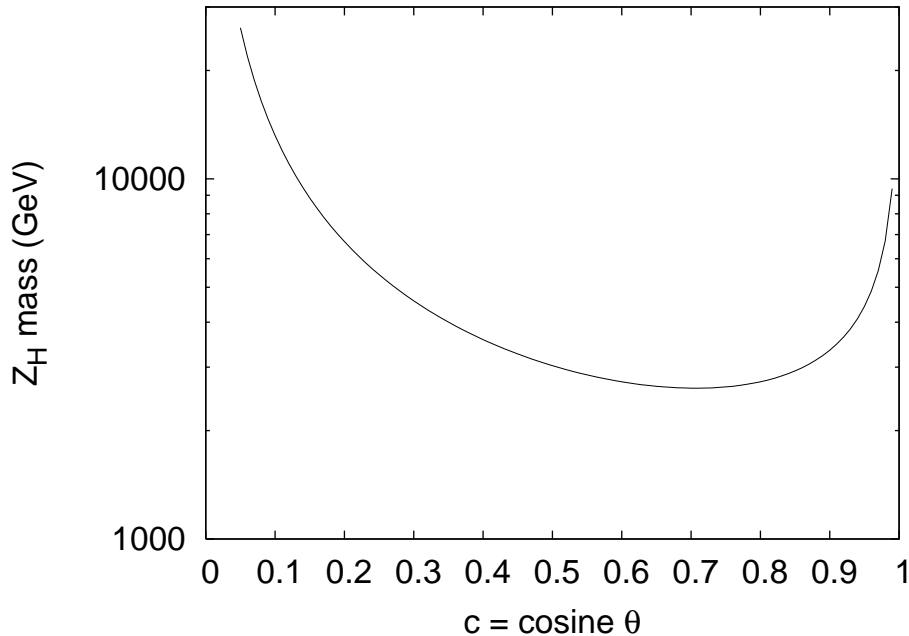


FIG. 1: The mass of the Z_H , for $f = 4000$ GeV.

and $y_e = \frac{3}{5}$ [3]. In this work, for simplicity, we always make the anomaly-cancelling choice.

The new particle content of the Littlest Higgs model is listed in Table I. All new particles have masses of order f , but not all are of equal importance of an experimental test of this model or one similar to it. The presence and couplings of the new $SU(2)$ gauge bosons Z_H, W_H^\pm would provide essential tests of the divergence-canceling mechanism of the model. The new vectorlike fermion T is necessary to cancel divergences from the top quark, but is not structurally related to the rest of the Littlest Higgs model, and so would not have great diagnostic power. This feature is specific to the Littlest Higgs model; in the Simplest Higgs model [7], for example, the new fermion content enters in a less *ad hoc* way.

The A_H is necessary to offset the divergence due to the loop contribution from Standard Model hypercharge, but this is numerically negligible, and some variants of the Littlest Higgs model have omitted this particle altogether by not gauging the corresponding $U(1)$ group, leaving instead an ‘unconsumed’ pseudoscalar[8]. In addition, the properties of the A_H vary greatly depending on the allocation of hypercharge across the two $U(1)$ groups[9]. The couplings of the Z_H are, in contrast, highly constrained and central to the structure of the Littlest Higgs model.

The mass of the Z_H depends on the overall symmetry-breaking scale f and the $SU(2)$

f	$SU(5)/SO(5)$ symmetry-breaking vev	4.0 TeV
c	Cosine of the $SU(2)_1 \otimes SU(2)_2$ mixing angle	0.05 to 0.99
c'	Cosine of the $U(1)_1 \otimes U(1)_2$ mixing angle	0.86
λ_1	Top sector Yukawa parameter	1.0
y_u	Quark hypercharge assignment parameter	$-\frac{2}{5}$
y_e	Lepton hypercharge assignment parameter	$+\frac{3}{5}$
v'	Scalar triplet vev	1.0 GeV

TABLE II: Free parameters of the Littlest Higgs model [3], with the values used in this study.

mixing angle, as displayed in Figure 1. In this work f is set to 4 TeV, largely to avoid direct-search exclusion limits already established at the Tevatron[9]. At this scale the mass ranges from a minimum of 2.6 TeV to in excess of 20 TeV, depending on the mixing angle. Since the Large Hadron Collider search reach for massive neutral gauge bosons is limited to roughly 5 TeV[10], the Z_H would not be experimentally accessible for $c \lesssim 0.27$ at $f = 4$ TeV. The search for the Z_H would in that case have to await an upgraded LHC or a future, even higher energy collider.

The free parameters of the Littlest Higgs model and the values used in this study are summarized in Table II.

In Section II of this paper, the two-body decays of the Z_H are described. Section III presents three-body decays, compares them to two-body decays, and identifies a symptom of perturbative breakdown. In Section IV, we summarize our results and draw some conclusions.

II. TWO-BODY DECAYS

The Z_H has a large variety of decays to two particles in the final state. In a reliable perturbative expansion we expect these to be the dominant decay channels, and so to collectively approximate the total decay width. There are decays to all Standard Model fermion-antifermion pairs with partial widths:

$$\Gamma(Z_H \rightarrow f\bar{f}) = \frac{N_c m_{Z_H}}{6\pi} \left(1 - 4\frac{m_f^2}{m_{Z_H}^2}\right)^{\frac{1}{2}} \left(\frac{gc}{4s}\right)^2 \left(1 - \frac{m_f^2}{m_{Z_H}^2}\right) \quad (3)$$

where N_c is the number of fermion colors, m_{Z_H} is the mass of the Z_H , m_f is the mass of the Standard Model fermion¹, c and s are the cosine and sine of the mixing angle θ diagonalizing the $[SU(2)]^2$ gauge bosons, and g is the Standard Model weak coupling constant. There are decays to standard model W^+W^- pairs with partial width:

$$\Gamma(Z_H \rightarrow W^+W^-) = g_{Z_H W_L W_L}^2 \frac{m_{Z_H}}{192\pi} \left(\frac{m_{Z_H}}{m_W} \right)^4 \left(1 - 4 \frac{m_W^2}{m_{Z_H}^2} \right)^{\frac{3}{2}} \left(1 + 20 \frac{m_W^2}{m_{Z_H}^2} + 12 \frac{m_W^4}{m_{Z_H}^4} \right) \quad (4)$$

where m_W is the mass of the W boson and the coupling constant is:

$$g_{Z_H W_L W_L} = \frac{g}{2} \frac{v^2}{f^2} s c (c^2 - s^2). \quad (5)$$

Decays to the Standard Model Z^0 boson and a Higgs boson occur with partial width:

$$\Gamma(Z_H \rightarrow Z^0 H) = \frac{|\vec{P}| V^2}{24\pi m_{Z_H}^2} \left(2 + \frac{E^2}{m^2} \right) \quad (6)$$

where $|\vec{P}|$, E , and m are the fixed momentum, energy, and mass, respectively, of the outgoing Z^0 in the Z_H rest frame, and the vertex factor constant V is:

$$V = \frac{g^2}{2c_w} \frac{c^2 - s^2}{2sc} v, \quad (7)$$

with c_w the cosine of the Standard Model weak mixing angle and v the electroweak scale $v = 246$ GeV. If the final-state bosons have negligible mass compared to the Z_H mass, the partial widths $\Gamma(Z_H \rightarrow W^+W^-)$ and $\Gamma(Z_H \rightarrow Z^0 H)$ reduce to equivalent expressions, as required by the Goldstone boson equivalence theorem, and as can be seen in Figure 2.

The Z_H also has decays to final states with one or more particles beyond those in the standard model. Since all the new particle masses in the Littlest Higgs model depend, at least weakly, on the free parameters of the theory, a given decay will typically be kinematically possible only for some portion of the parameter space. All the decay branches described here are possible for at least some region of parameter space. It is convenient to group these decays according to final state spins. First, decays to a vector boson plus a scalar take the same form as to $Z^0 H$:

$$\Gamma(Z_H \rightarrow \text{Vector} + \text{Scalar}) = \frac{|\vec{P}| V^2}{24\pi m_{Z_H}^2} \left(2 + \frac{E^2}{m^2} \right) \quad (8)$$

¹ In the numerical results presented below, all Standard Model fermion masses other than the top quark mass have been omitted since they are utterly negligible compared to the mass of the Z_H .

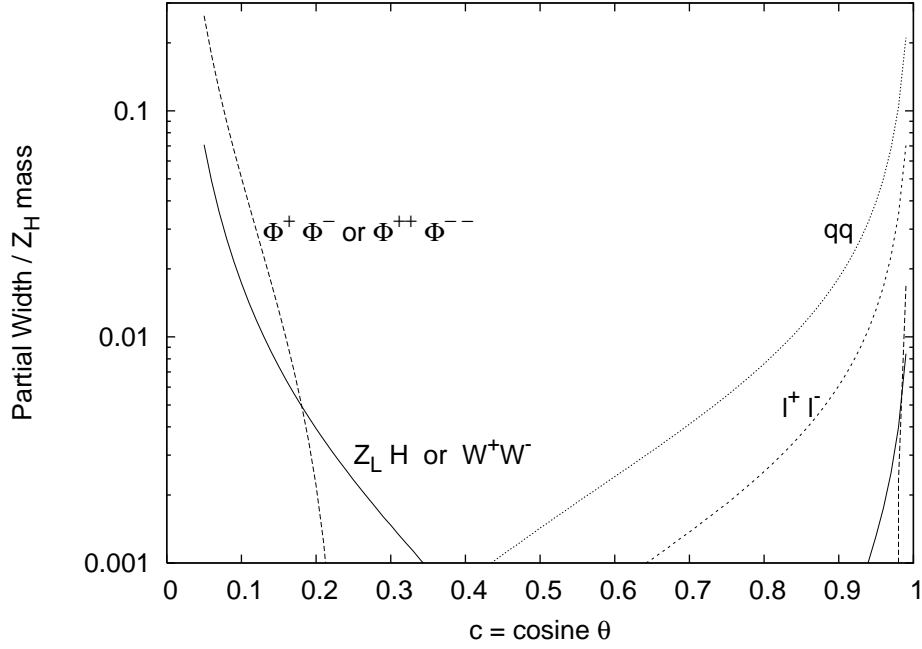


FIG. 2: The partial width to mass ratio of the Z_H for significant two-body decays, at $f = 4$ TeV. The curve labeled qq is to any Standard Model quark-antiquark pair (since even the top mass is negligible at this scale), and the curve l^+l^- is to electrons, and identically to any Standard Model lepton-antilepton pair.

where $|\vec{P}|$, E , and m are the momentum, energy, and mass of the outgoing vector boson. The kinematically allowed processes of this type are $Z_H \rightarrow Z^0\phi^0, A_H\phi^0, W^\pm\phi^\mp$ and $A_H H$. The vertex constant for each process is presented in [3].

Various decays to two scalars are possible. The general form of this partial width is:

$$\Gamma(Z_H \rightarrow \text{Scalar} + \text{Scalar}) = \frac{|\vec{P}|^3}{6\pi m_{Z_H}^2} V^2. \quad (9)$$

The possible decays of this type are $Z_H \rightarrow H\phi^P, \phi^0\phi^P, \phi^+\phi^-,$ and $\phi^{++}\phi^{--}$. Vertex constants V again are as in [3].

Finally, decays to top quark plus the new antifermion \bar{T} (and the conjugate process) are possible. The kinematics of this process are complicated by the differing masses of the final state fermions, with result:

$$\Gamma(Z_H \rightarrow t\bar{T}) = \frac{N_c |\vec{P}|}{6\pi} g_v^2 \left(1 - \frac{m_t^2}{m_{Z_H}^2} - \frac{m_T^2}{m_{Z_H}^2} + \left(1 - \frac{(m_T - m_t)^2}{m_{Z_H}^2} \right)^2 \right) \quad (10)$$

$Z_L H H$	All c
$Z_L W_L^+ W_L^-$	All c
$A_H H H$	All c
$A_L W_L^+ W_L^-$	All c
$A_H W_L^+ W_L^-$	All c
$W_L^+ H \phi^-$	$c \leq 0.49$ or $c \geq 0.88$
$Z_L H \phi^0$	$c \leq 0.48$ or $c \geq 0.88$
$A_H H \phi^0$	$c \leq 0.38$ or $c \geq 0.93$
$Z_L \phi^0 \phi^0$	$c \leq 0.23$ or $c \geq 0.98$
$Z_L \phi^p \phi^p$	$c \leq 0.23$ or $c \geq 0.98$
$Z_L \phi^{++} \phi^{--}$	$c \leq 0.23$ or $c \geq 0.98$
$A_L \phi^{++} \phi^{--}$	$c \leq 0.23$ or $c \geq 0.98$
$W_L^+ \phi^0 \phi^-$	$c \leq 0.23$ or $c \geq 0.98$
$W_L^+ \phi^p \phi^-$	$c \leq 0.23$ or $c \geq 0.98$
$W_L^+ \phi^+ \phi^{--}$	$c \leq 0.23$ or $c \geq 0.98$
$A_H \phi^0 \phi^0$	$c \leq 0.20$ or $c \geq 0.98$
$A_H \phi^p \phi^p$	$c \leq 0.20$ or $c \geq 0.98$
$A_H \phi^{++} \phi^{--}$	$c \leq 0.20$ or $c \geq 0.98$

TABLE III: The kinematically allowed ranges of the $SU(2)$ mixing angle c , for the allowed decays of the Littlest Higgs Z_H to three particle final states, at $f = 4$ TeV and $c' = 0.86$.

where g_v is the vector coupling (equal and opposite to the axial coupling) for these particles[3].

The partial widths of these various processes vary greatly with c , as shown in Figure 2. At the scale $f = 4$ TeV, even the effect of the top quark mass is negligible, so decays to top pairs have the same width as to other quark-antiquark pairs. Note that decays to massive Φ scalars dominate for $c \lesssim 0.2$, which has been omitted in some prior studies, perhaps due to concerns about perturbativity. All two-body processes not included in the figure contribute a partial width of less than one percent of the mass for all values of c .

Although Standard Model final state masses are negligible at $f = 4$ TeV, much lower symmetry-breaking scales are not excluded by direct searches for some regions of parameter

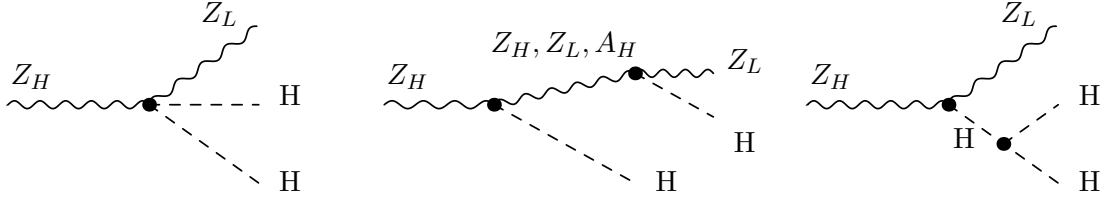


FIG. 3: The Feynman diagrams contributing to the decay process $Z_H \rightarrow Z_L HH$.

space[9]. At $f = 1.5$ TeV, the final state mass effects are on the order of ten percent.

III. THREE-BODY DECAYS

We surveyed the available Littlest Higgs couplings and particle masses to determine the possible three-body decays of the Z_H . There are eighteen distinct kinematically-allowed three-body tree-level decays for the Z_H , which are listed in Table III. Many of these decay processes are possible only for small or very large values of c (where Z_H becomes very massive) due to the presence of one or more TeV-scale particles in the final state. Note that no decays to the W_H are possible because it has the same mass as the Z_H to high order in $\frac{v}{f}$.

These decays will characteristically each involve a set of eight or so Feynman diagrams, including one with a quadrilinear vertex factor. We exclude from our list processes without a quadrilinear coupling and with a related kinematically possible two-body decay, since these would constitute a double-counting of a two-body decay with an added decay step for one of the decay products.

The diagrams for the process $Z_H \rightarrow Z_L HH$ are displayed in Figure 3. Note that a diagram with a ϕ particle propagator is not present because the combined effects of a factor of $\frac{v}{f}$ at each vertex places this diagram at higher order in the sigma expansion than we are working.

Since the calculation of a three-body decay is analytically difficult, especially because we wish to retain final state masses, we implemented a Monte Carlo integration to determine the partial width for $Z_H \rightarrow Z_L HH$. This partial width becomes very large at large and

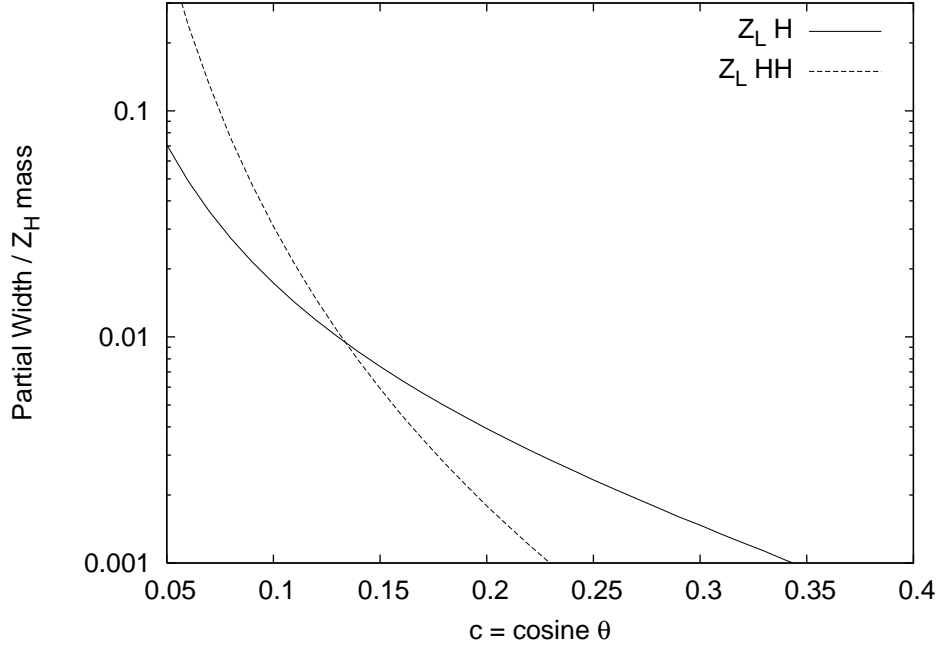


FIG. 4: The partial width to mass ratio of the Z_H decays to the $Z_L H$ and $Z_L HH$ final states, at $f = 4$ TeV. The decay width to $Z_L HH$ exceeds that to $Z_L H$ for $c < 0.13$, signalling a breakdown in perturbative reliability. In the approach to this region, increasing decay widths to multiple scalars suggest a distinctive collider signature.

especially at small c . We would expect this behavior because many vertex factors contain both c and s in the denominator. In fact this effect is present in two-body decays, but for three-body decays the width divergence is much more severe. The width-to-mass ratios for the processes $Z_H \rightarrow Z_L H$ and $Z_H \rightarrow Z_L HH$ are shown in Figure 4. The three-body partial width exceeds the two-body partial width for $c < 0.13$.

This is clearly pathological, since three-body decays should be coupling-suppressed compared to two-body decays. What is happening? The scalar-gauge boson couplings originate in the scalar kinematic Lagrangian term[3]:

$$\mathcal{L}_\Sigma = \frac{f^2}{8} (\mathcal{D}_\mu \Sigma) (\mathcal{D}^\mu \Sigma)^\dagger \quad (11)$$

where \mathcal{D}_μ is the gauge-covariant derivative in the bifundamental representation of $SU(5)$,

$$\mathcal{D}_\mu = \partial_\mu \Sigma - i \sum_{j=1}^2 \left(g_j (W_j \Sigma + \Sigma W_j^T) + g'_j (B_j \Sigma + \Sigma B_j^T) \right) \quad (12)$$

and Σ is the nonlinear representation:

$$\Sigma = e^{2i\Pi/f}\Sigma_0 \quad (13)$$

and where Π is a 5 x 5 matrix which parametrizes the scalars in the theory, including the longitudinal modes of massive Standard Model gauge bosons as well as the new TeV scale gauge bosons. The constant matrix Σ_0 is the choice of vacuum expectation value for this model[2]. In the gauge-covariant derivative the index j is over the pairs of electroweak gauge groups, and the index over the generators of $SU(2)$ are suppressed.

As Σ is expanded in a power series, an infinite set of vertex factors will result from the scalar kinematic term, each involving two gauge bosons, some number of scalars and a power of $\frac{v}{f_{sc}}$, with the factor $\frac{1}{sc}$ appearing due to mixing of the two $SU(2)$ gauge groups. The lowest order terms will be two gauge bosons plus one scalar, then plus two scalars, then three scalars, and so on. The existence of vertex factors connecting more than four external particles signals the fact that any nonlinear sigma model is a nonrenormalizable effective theory. The effects of this nonrenormalizability would become evident in loop calculations, but here we are considering a different kind of divergence - that of the expansion of the Σ field itself. We might expect this divergence to become parametrically significant at about $\frac{v}{f_{sc}} \sim 1$, which at $f = 4$ TeV occurs at $c = 0.06$ [11]. Instead, we see the pathological effect of the decay to $Z_L H H$ exceeding the decay to $Z_L H$ at as high as $c = 0.13$, as we approach the strong coupling regime.

We have not presented the partial widths for other three-body processes, but can note that the decay $Z_H \rightarrow Z_L W_L^+ W_L^-$ should approximate that of $Z_L H H$ due to the Goldstone boson equivalence across these processes. With the exception of $Z_H \rightarrow A_L W_L^+ W_L^-$, all other processes will be kinematically suppressed due to the presence of one, two, or three TeV-scale particles in the final state. These processes will in fact be kinematically excluded except for small or very large c . By combining these contributions we can roughly estimate that the total partial width to three final state particles is less than an order of magnitude larger than that to $Z_L H H$ for all but the smallest and largest values of c .

IV. CONCLUSIONS

We have found that a characteristic three-body decay can become large at small and very large values of the $SU(2)$ mixing angle c , and exceeds a related two-body decay at $c < 0.13$, indicating divergence of the sigma field expansion. The pathology of the Littlest Higgs model in this region of parameter space is due to a nonlinear sigma scalar field representation, and occurs over a larger regions of parameter space than a simple parametric calculation would suggest.

Although perturbative calculations break down in this region, we can note an interesting qualitative feature: The Littlest Higgs model near the nonperturbative region would produce decays with high scalar multiplicities, especially with multiple Higgs boson production, and so with large b quark multiplicities. This feature of large rates of scalar production in decays may be a general feature of nonlinear sigma fields near their perturbative limits, as the regime of strong dynamics is approached. In particular, we would expect similar results for decays of the Littlest Higgs A_H and W_H^\pm , and for the new gauge bosons of other models in the Little Higgs class, such as the Simplest Higgs Z' . We should note, however, that all results presented here are at tree-level only, and could be significantly modified by the effects of loop corrections.

We have presented the various expressions for two-body decays of the Littlest Higgs Z_H , including all final state masses. The numerical width-to-mass ratios for the most significant of these processes, and for a characteristic three-body decay process, confirm that the Z_H is a narrow resonance, as needed for dilepton invariant mass searches [12].

Acknowledgment

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